

Tension in the inclusive versus exclusive determinations of $|V_{cb}|$: A possible role of new physics

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We reconsider the possibility that the tension in the $|V_{cb}|$ determinations from inclusive and exclusive B decay modes is due to a new physics effect. We modify the Standard Model effective Hamiltonian for semileptonic $b \rightarrow c$ transitions including a tensor operator with a lepton flavour dependent coupling ϵ_T^ℓ , and investigate separately the muon and electron modes. The interference term between SM and NP, proportional to the lepton mass, has different impact in the inclusive and exclusive B modes to muon. Moreover, even when the lepton mass is small as for the electron, the NP effect is different in inclusive and exclusive B channels. For both μ and e we find a region of $\epsilon_T^{\mu,e}$ where the constraints from $B^- \rightarrow D^{(*)0} \ell^- \bar{\nu}_\ell$ and $B \rightarrow X_c \ell \bar{\nu}_\ell$ are satisfied for the same $|V_{cb}|$.

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Introduction. The precise determination of the Standard Model (SM) parameters is a fundamental step towards the search for new physics (NP). In particular, the CKM matrix elements play a peculiar role, being related to the SM description of CP violation in the quark sector. Special cases are V_{ub} and V_{cb} , whose ratio $|V_{ub}|/|V_{cb}|$ appears in the length of one side of the unitarity triangle. They can be measured in several ways, in particular using semileptonic decays of b -flavoured hadrons, and for both $|V_{ub}|$ and $|V_{cb}|$ the results obtained from inclusive and exclusive B semileptonic modes are only marginally compatible. From exclusive decays, the average in [1]

$$|V_{cb}|_{excl} = (39.78 \pm 0.42) \times 10^{-3} \quad (1)$$

is based on experimental data and on hadronic quantities computed by lattice QCD [2–5]. The Particle Data Group is more conservative: $|V_{cb}|_{excl} = (39.2 \pm 0.7) \times 10^{-3}$ [6], while in a recent study of the single $B \rightarrow D$ mode $|V_{cb}| = (40.49 \pm 0.97) \times 10^{-3}$ is obtained [7].

Results from inclusive B decays are summarized in [6, 9]:

$$|V_{cb}|_{incl} = (42.21 \pm 0.78) \times 10^{-3}. \quad (2)$$

The tension between (1) and (2), one of the so-called *flavour anomalies* [10], affects the SM predictions for several other observables, as discussed, e.g., in [11].

Another flavour anomaly are the ratios $\mathcal{R}(D^{(*)}) = \frac{\mathcal{B}(B \rightarrow D^{(*)} \tau \bar{\nu}_\tau)}{\mathcal{B}(B \rightarrow D^{(*)} \mu \bar{\nu}_\mu)}$, that exceed the SM prediction [12–14].¹ Surprisingly, both the $|V_{cb}|$ and $\mathcal{R}(D^{(*)})$ anomalies involve tree-level processes, and the second one points to violation of lepton flavour universality (LFU). In [16] it has been shown that an additional tensor operator in the effective Hamiltonian inducing semileptonic $b \rightarrow c$ decays could accommodate the experimental findings for $\mathcal{R}(D^{(*)})$, assuming that it contributes only for τ leptons.

The coupling ϵ_T^τ of the new term in the effective Hamiltonian is constrained by the $R(D^{(*)})$ data in a region allowing to reconcile the theoretical result with measurement. Hence, it is worth scrutinizing if this kind of violation of LFU is at work also in the case of the $|V_{cb}|$ anomaly.

The possibility that new physics is involved in the difference between inclusive and exclusive determinations of $|V_{cb}|$ has been previously considered [17–19]. It has been excluded using the argument that a new scalar or a new tensor operator in the effective $b \rightarrow c$ Hamiltonian, for massless leptons, produces the same effect in both exclusive and inclusive semileptonic modes at zero recoil, inducing the same changes in $|V_{cb}|$. Moreover, introducing a new vector or axial-vector structure in the Hamiltonian leads to modified W couplings, and this, due to the $SU(2)$ symmetry of SM, in turn modifies the Z couplings to fermions at a level experimentally excluded [18].

However, in the inclusive B mode the determination of $|V_{cb}|$ goes through the measurement of the moments of the full lepton energy spectrum and of the hadronic mass distribution, to obtain the OPE parameters in the theoretical expression of the inclusive width. $|V_{cb}|$ is then determined comparing the theoretical and experimental width [20]. This differs from the exclusive modes that are analyzed close to zero recoil [6]. As we show for the new tensor operator, the lepton mass, in the case of muons, can be large enough to produce a sizable interference between the SM and the NP contribution, with different effects in the exclusive and in the inclusive B decays and a different impact on $|V_{cb}|$, jeopardizing the argument in [18]. For this reason, whenever possible we consider separately the muon and electron modes. Moreover, we find that the NP effect is not the same in the inclusive and exclusive channels, and this also in the electron case where the interference between SM and NP is negligible.

To be specific, we extend the effective Hamiltonian governing the $b \rightarrow c \ell \bar{\nu}_\ell$ transitions as follows [16, 21, 22]:

$$H_{eff} = \frac{G_F}{\sqrt{2}} V_{cb} \left[\bar{c} \gamma_\mu (1 - \gamma_5) b \bar{\ell} \gamma^\mu (1 - \gamma_5) \bar{\nu}_\ell + \epsilon_T^\ell \bar{c} \sigma_{\mu\nu} (1 - \gamma_5) b \bar{\ell} \sigma^{\mu\nu} (1 - \gamma_5) \bar{\nu}_\ell \right]. \quad (3)$$

¹ A recent analysis with references to previous studies is in [15].

We assume that not only ϵ_T^τ [16], but also $\epsilon_T^{(\mu,e)}$ can be different from zero, and we investigate the effect of the new term for light leptons and the consequences for $|V_{cb}|$.

Inclusive $B \rightarrow X_c \ell \bar{\nu}_\ell$ decays. The heavy quark expansion (HQE) [23–26] allows to write the inclusive decay width of heavy hadrons (H_Q) as a series in powers of the inverse heavy quark mass. Each term is the product of coefficient functions times the $\langle H_Q | O_i | H_Q \rangle$ matrix elements of local operators O_i of increasing dimension; at each order in $1/m_Q$, the coefficient functions can be further expanded in α_s . The leading term corresponds to the free Q decay width; the $\mathcal{O}(m_Q^{-1})$ term is absent. Specializing to B meson, the expression of $B \rightarrow X_c \ell \bar{\nu}_\ell$ decay width in SM for massless leptons, with $\mathcal{O}(\alpha_s^2)$ and $\mathcal{O}(1/m_b^3)$ corrections, can be found in [8]. Here we compute $\Gamma(B \rightarrow X_c \ell \bar{\nu}_\ell)$ including the contribution of the tensor operator in Eq. (3) and considering massive leptons. The decay distribution in the dilepton invariant mass $\hat{q}^2 = q^2/m_b^2$ reads

$$\frac{d\Gamma}{d\hat{q}^2} = C(q^2) \left[\frac{d\tilde{\Gamma}}{d\hat{q}^2}|_{SM} + |\epsilon_T|^2 \frac{d\tilde{\Gamma}}{d\hat{q}^2}|_{NP} + \text{Re}(\epsilon_T) \frac{d\tilde{\Gamma}}{d\hat{q}^2}|_{INT} \right], \quad (4)$$

with $C(q^2) = \frac{G_F^2 |V_{cb}|^2 m_b^5}{96\pi^3} \lambda^{1/2} \left(1 - \frac{\hat{m}_\ell^2}{\hat{q}^2}\right)^2$, $\hat{m}_\ell = m_\ell/m_b$, $\lambda = \lambda(1, \rho, \hat{q}^2)$ the triangular function and $\rho = m_c^2/m_b^2$.

Using the HQE, each term $A=SM, NP, INT$ in (4) can be written as

$$\frac{d\tilde{\Gamma}}{d\hat{q}^2}|_A = \frac{d\tilde{\Gamma}_0}{d\hat{q}^2}|_A - \frac{\mu_\pi^2 - \mu_G^2}{2m_b^2} \frac{d\tilde{\Gamma}_{1/m_b^2}^{(1)}}{d\hat{q}^2}|_A + \frac{\mu_G^2}{m_b^2} \frac{d\tilde{\Gamma}_{1/m_b^2}^{(2)}}{d\hat{q}^2}|_A. \quad (5)$$

The leading order terms in the $1/m_b$ expansion (5) read:

$$\begin{aligned} \frac{d\tilde{\Gamma}_0}{d\hat{q}^2}|_{SM} &= (1-\rho)^2 \left(1 + 2\frac{\hat{m}_\ell^2}{\hat{q}^2}\right) + \\ &\quad (1+\rho)\hat{q}^2 \left(1 - \frac{\hat{m}_\ell^2}{\hat{q}^2}\right) - \hat{q}^4 \left(2 + \frac{\hat{m}_\ell^2}{\hat{q}^2}\right), \\ \frac{d\tilde{\Gamma}_0}{d\hat{q}^2}|_{NP} &= 8 \left(1 + 2\frac{\hat{m}_\ell^2}{\hat{q}^2}\right) [2(1-\rho)^2 - \hat{q}^2(1+\rho+\hat{q}^2)], \\ \frac{d\tilde{\Gamma}_0}{d\hat{q}^2}|_{INT} &= -36\sqrt{\rho}\hat{m}_\ell(1-\rho+\hat{q}^2). \end{aligned} \quad (6)$$

For the $1/m_b^2$ corrections we find:

$$\frac{d\tilde{\Gamma}_{1/m_b^2}^{(1)}}{d\hat{q}^2}|_A = \frac{d\tilde{\Gamma}_0}{d\hat{q}^2}|_A \quad (A = SM, NP, INT), \quad (7)$$

$$\begin{aligned} \frac{d\tilde{\Gamma}_{1/m_b^2}^{(2)}}{d\hat{q}^2}|_{SM} &= 2 \left\{ -\lambda \left(2 + \frac{\hat{m}_\ell^2}{\hat{q}^2}\right) - [-3 + 8\rho - 3\rho^2 + 3\hat{q}^2(1+\rho)] \left(1 + \frac{\hat{m}_\ell^2}{\hat{q}^2}\right) - \frac{\hat{m}_\ell^2}{\hat{q}^2}(3-\rho) \right. \\ &\quad \left. + \frac{1}{\lambda} \left[(1-\rho)^2 \left[3\rho \left(1 + \frac{\hat{m}_\ell^2}{\hat{q}^2}\right) - \left(1 - \frac{\hat{m}_\ell^2}{\hat{q}^2}\right) \right] + \hat{q}^2 \left[2(1+3\rho) - (1+12\rho+3\rho^2) \left(1 + \frac{\hat{m}_\ell^2}{\hat{q}^2}\right) \right] \right] \right\}, \\ \frac{d\tilde{\Gamma}_{1/m_b^2}^{(2)}}{d\hat{q}^2}|_{NP} &= -16 \left(1 + 2\frac{\hat{m}_\ell^2}{\hat{q}^2}\right) \left\{ \lambda + \rho(5-3\rho) + 3\hat{q}^2(1+\rho) + \frac{1}{\lambda} [\hat{q}^2(-1+6\rho+3\rho^2) - (1-\rho)^2(-1+3\rho)] \right\}, \\ \frac{d\tilde{\Gamma}_{1/m_b^2}^{(2)}}{d\hat{q}^2}|_{INT} &= 24\sqrt{\rho}\hat{m}_\ell \left\{ -2 - 3\hat{q}^2 + 3\rho + \frac{1}{\lambda} [3(1-\rho)^2 - \hat{q}^2(1+3\rho)] \right\}. \end{aligned} \quad (8)$$

μ_π^2 and μ_G^2 parameterize the matrix elements $\langle B | O_i | B \rangle$ of dimension 5 operators. Setting $\epsilon_T = 0$ we recover the SM result for massive leptons [27]. In our numerical analysis we also include $1/m_b^3$ and $\mathcal{O}(\alpha_s^2)$ corrections in the SM contribution, neglecting the lepton mass [8]. We do not include such corrections in the NP term, which has to be small with respect to the SM one. Hence, we write the inclusive semileptonic width as follows:

$$\begin{aligned} \Gamma_{sl} &= \left(\int_{\hat{q}_{min}^2}^{\hat{q}_{max}^2} \frac{d\Gamma}{d\hat{q}^2} d\hat{q}^2 \right) + \Gamma_0 \left[a^{(1)} \frac{\alpha_s(m_b)}{\pi} \right. \\ &\quad \left. + a^{(2,\beta_0)} \beta_0 \left(\frac{\alpha_s}{\pi} \right)^2 + a^{(2)} \left(\frac{\alpha_s}{\pi} \right)^2 + p^{(1)} \frac{\alpha_s}{\pi} \frac{\mu_\pi^2}{m_b^2} \right] \end{aligned}$$

$$+ g^{(1)} \frac{\alpha_s}{\pi} \frac{\mu_G^2}{m_b^2} + d^{(0)} \frac{\rho_D^3}{m_b^3} - g^{(0)} \frac{\rho_{LS}^3}{m_b^3} \Big]. \quad (9)$$

The first integrand is in Eq.(4). Γ_0 is the LO SM width, ρ_D^3 , ρ_{LS}^3 parameterize dimension 6 operator matrix elements, and the various coefficients are reported in [8]. We use their values in the kinetic scheme, while for m_c we use the \overline{MS} result at the scale $\mu = 3$ GeV, as in [28]. The Particle Data Group quotes [6]:

$$\mathcal{B}(B^+ \rightarrow X_c e^+ \nu_e) = (10.8 \pm 0.4) \times 10^{-2}. \quad (10)$$

Since no data are separately reported for muon, we use (10) also in that case. In Fig. 1 we show the result of exploiting the datum Eq. (10) at 1σ level to constrain

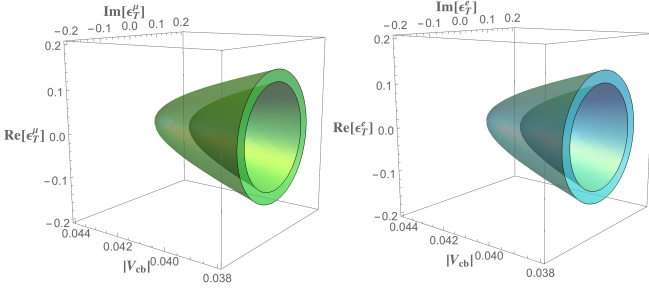


FIG. 1: Cutaway view of the parameter space $\text{Re}(\epsilon_T^\ell)$, $\text{Im}(\epsilon_T^\ell)$, and $|V_{cb}|$ allowing to obtain $\mathcal{B}(B^+ \rightarrow X_c \ell^+ \nu_\ell)$ within 1σ for muon (left) and electron mode (right).

$(\text{Re}(\epsilon_T^\ell), \text{Im}(\epsilon_T^\ell), |V_{cb}|)$ for the muon and electron mode. We keep the lepton masses distinct in the two cases. The inclusive branching ratio bounds $|V_{cb}|$ from above: imposing the 1σ constraint, one finds $|V_{cb}| \leq 42.85 \times 10^{-3}$ in the muon case, and $|V_{cb}| \leq 42.73 \times 10^{-3}$ in the electron case, in correspondence to the SM point $\epsilon_T^\ell = 0$. When ϵ_T^ℓ is allowed to deviate from zero, the hollow regions in Fig. 1 are found, which continue to smaller values of $|V_{cb}|$ if $|\epsilon_T^\ell|$ is increased; however, a lower bound on $|V_{cb}|$ is set by the exclusive modes.

Exclusive $B \rightarrow D^{()}$ modes.* The theoretical description of $B \rightarrow D^{(*)} \ell \bar{\nu}_\ell$ modes requires the $B \rightarrow D^{(*)}$ form factors. In the HQ limit all such form factors can be expressed in terms of the Isgur-Wise function $\xi(w)$ [29], with w related to the dilepton invariant mass, $q^2 = m_B^2 + m_{D^{(*)}}^2 - 2m_B m_{D^{(*)}} w$. Corrections involve both $1/m_Q$ and $\mathcal{O}(\alpha_s)$ terms, and can be found, e.g., in [30, 31]. For finite quark mass, several form factor determinations are available. For $F_1(q^2)$ and $F_0(q^2)$ in the $B \rightarrow D$ matrix element of the weak current $j_\mu = \bar{c} \gamma_\mu (1 - \gamma_5) b$, the lattice QCD results in [3] have been used to obtain the average Eq. (1). For consistency, we use them. Moreover, we determine the form factor $F_T(q^2)$ parametrizing the matrix element of the tensor current, $j_{\mu\nu} = \bar{c} \sigma_{\mu\nu} (1 - \gamma_5) b$, from $F_1(q^2)$ and the HQ relation at NLO [30, 31], following the procedure described in details in [16]. The standard parameterization of the $B \rightarrow D^*$ matrix element of the currents j_μ and $j_{\mu\nu}$ involves the form factors V , A_i ($i = 0, 3$), and T_i ($i = 0, 5$), as defined in [16]. Lattice QCD results are only available for A_1 at the zero recoil $w = 1$ [5]: we determine the other ones using again HQ relations at NLO [16].

Starting from H_{eff} in (3), the differential $B \rightarrow M_c \ell \bar{\nu}$ decay distribution, with $M_c = D, D^*$ can be written as

$$\frac{d\Gamma}{dq^2}(B \rightarrow M_c \ell \bar{\nu}_\ell) = \frac{d\Gamma}{dq^2}\Big|_{SM} + \frac{d\Gamma}{dq^2}\Big|_{NP} + \frac{d\Gamma}{dq^2}\Big|_{INT}, \quad (11)$$

with the three terms for massive leptons given in [16].

For the modes $B^- \rightarrow D^0 \ell^- \bar{\nu}_\ell$, the BaBar Collaboration has reported separate measurements for μ and e [4]:

$$\mathcal{B}(B^- \rightarrow D^0 \mu^- \bar{\nu}_\mu) = (2.25 \pm 0.04 \pm 0.17) \times 10^{-2} \quad (12)$$

$$\mathcal{B}(B^- \rightarrow D^0 e^- \bar{\nu}_e) = (2.38 \pm 0.04 \pm 0.15) \times 10^{-2}. \quad (13)$$

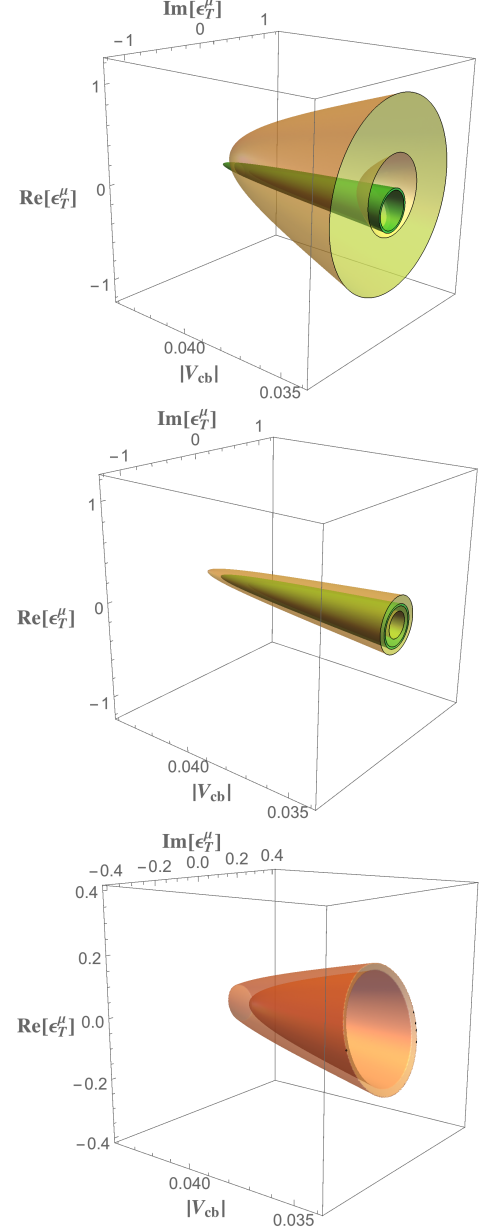


FIG. 2: Muon modes: allowed regions in the $(\text{Re}(\epsilon_T^\mu), \text{Im}(\epsilon_T^\mu), |V_{cb}|)$ parameter space, determined from the inclusive mode (green hollow region) together with the exclusive D mode (yellow hollow region, upper plot), and from the inclusive mode together with the decay to D^* (yellow hollow region, middle plot). The intersection of the three regions is shown in the lower plot.

We constrain $(\text{Re}(\epsilon_T^\ell), \text{Im}(\epsilon_T^\ell), |V_{cb}|)$ comparing these data with the theoretical expressions including the errors on the form factors quoted in [3], and the uncertainty on the parameter $\bar{\Lambda} = m_{H_Q} - m_Q$ entering in the HQ relation between F_1 and F_T . We use the conservative value $\bar{\Lambda} = 0.5 \pm 0.2$ GeV.

In the case of $B^- \rightarrow D^{*0} \ell^- \bar{\nu}_\ell$, we consider the distribution $\frac{d\Gamma}{dw}$ and compare the theory prediction to exper-

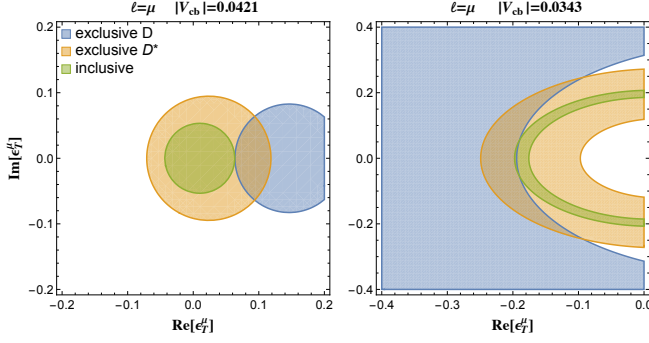


FIG. 3: Muon channel: projections of the overlap parameter region in Fig. 2 on the $(\text{Re}(\epsilon_T^\mu), \text{Im}(\epsilon_T^\mu))$ plane. In each panel, the blue region corresponds to the constraint from the exclusive decay to D , the orange region to the constraint from D^* , and the green region to the constraint from the inclusive mode. The left panel corresponds to the largest allowed value of $|V_{cb}|$, the right panel to the smallest allowed value of $|V_{cb}|$. Outside this range of $|V_{cb}|$, the parameter regions bound from the three decay modes do not overlap.

iment close to the zero recoil point $w \rightarrow 1$. The BaBar Collaboration fits the measured data using the expression

$$\begin{aligned} \frac{d\Gamma}{dw}(B \rightarrow D^* \ell \bar{\nu}_\ell) &= \frac{G_F^2 |V_{cb}|^2 m_B^5}{48\pi^3} (1 - r^*)^2 r^{*3} W_{D^*}(w) \\ h_{A_1}^2(w) \sqrt{w^2 - 1} (w + 1)^2 &\left\{ \left[1 + (1 - R_2(w)) \frac{w - 1}{1 - r^*} \right]^2 \right. \\ &\left. + 2 \left[\frac{1 - 2wr^* + r^{*2}}{(1 - r^*)^2} \right] \left[1 + R_1(w)^2 \frac{w - 1}{w + 1} \right] \right\} \quad (14) \end{aligned}$$

and $r^* = m_{D^*}/m_B$. The function $W_{D^*}(w)$, defined in [4], satisfies the condition $W_{D^*}(1) = (1 - \frac{m_\ell^2}{m_B^2(1-r^*)^2})^2 (1 + \frac{m_\ell^2}{2m_B^2(1-r^*)^2})$. For the three functions in (14) the parametrization is used [31]:

$$\begin{aligned} h_{A_1}(w) &= h_{A_1}(1) [1 - 8\hat{\rho}^2 z + (53\hat{\rho}^2 - 15)z^2 \\ &\quad - (231\hat{\rho}^2 - 91)z^3], \\ R_1(w) &= R_1(1) - 0.12(w - 1) + 0.05(w - 1)^2, \\ R_2(w) &= R_1(1) + 0.11(w - 1) - 0.06(w - 1)^2, \end{aligned} \quad (15)$$

with $z = \frac{\sqrt{w+1} - \sqrt{2}}{\sqrt{w+1} + \sqrt{2}}$. In [4] $\hat{\rho}^2$, $R_1(1)$, $R_2(1)$ are fitted separately for $\ell = \mu$ and $\ell = e$. Their values, together with the measurements $\mathcal{B}(B \rightarrow D^* \mu \bar{\nu}_\mu) = (5.34 \pm 0.06 \pm 0.37)\%$, $\mathcal{B}(B \rightarrow D^* e \bar{\nu}_e) = (5.50 \pm 0.05 \pm 0.23)\%$, give

$$h_{A_1}^\mu(1)|V_{cb}| = (35.63 \pm 1.96) \times 10^{-3} \quad (16)$$

$$h_{A_1}^e(1)|V_{cb}| = (35.94 \pm 1.65) \times 10^{-3}. \quad (17)$$

To bound $(\text{Re}(\epsilon_T^\ell), \text{Im}(\epsilon_T^\ell), |V_{cb}|)$, we compare the theoretical expression of $d\Gamma/dw$ for $w \rightarrow 1$,

$$\frac{d\Gamma^{th}}{dw}(B^- \rightarrow D^{*0} \ell^- \bar{\nu}_\ell)|_{w \rightarrow 1}$$

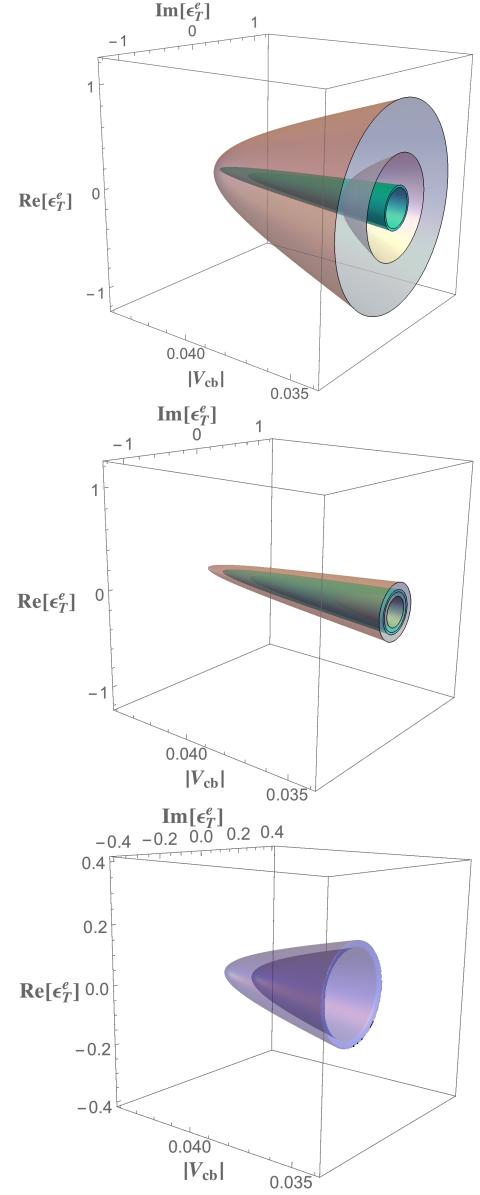


FIG. 4: Electron modes: allowed regions in the $(\text{Re}(\epsilon_T^e), \text{Im}(\epsilon_T^e), |V_{cb}|)$ parameter space from the inclusive mode (cyan hollow region) together with the exclusive decay to D (gray hollow region, upper plot), and from the inclusive mode together with D^* (gray hollow region, middle plot). The lower plot shows the intersection of the three regions.

$$\begin{aligned} &= \frac{G_F^2 |V_{cb}|^2 m_{D^*}^2}{16\sqrt{2}\pi^3 m_B} \sqrt{w - 1} \left[1 - \frac{m_\ell^2}{(m_B - m_{D^*})^2} \right]^2 \times \\ &\quad \left\{ (m_B + m_{D^*})^2 [2(m_B - m_{D^*})^2 + m_\ell^2] A_1(1)^2 \right. \\ &\quad + |\epsilon_T|^2 4[(m_B - m_{D^*})^2 + 2m_\ell^2] [m_B \tilde{T}_1(1) + m_{D^*} \tilde{T}_2(1)]^2 \\ &\quad \left. - 12 \text{Re}(\epsilon_T) (m_B^2 - m_{D^*}^2) m_\ell A_1(1) [m_B \tilde{T}_1(1) + m_{D^*} \tilde{T}_2(1)] \right\}, \quad (18) \end{aligned}$$

to Eq. (14) for $w \rightarrow 1$, using (16) and (17) for the muon and the electron mode, respectively. In (18), \tilde{T}_i are combinations of the tensor form factors $\tilde{T}_0 = T_0 - T_5$,

$$\tilde{T}_1 = T_1 + T_3 \text{ and } \tilde{T}_2 = T_2 + T_4.$$

We are now able to put together the constraints from the B inclusive and exclusive D, D^* decay modes. For

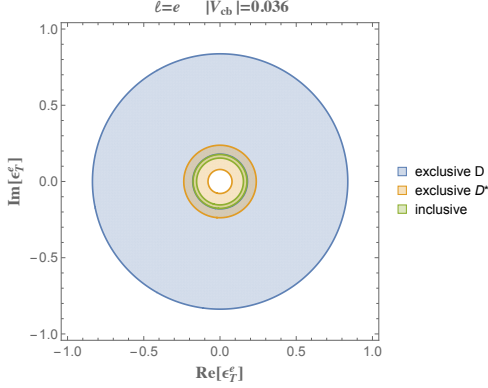


FIG. 5: Electron channel: projections of the overlap region in Fig. 4 in the $(\text{Re}(\epsilon_T^e), \text{Im}(\epsilon_T^e))$ plane. $|V_{cb}|$ is set to the smallest allowed value. Colour coding as in Fig. 3.

the muon channel, we show in the two upper panels of Fig. 2 the parameter regions selected by the exclusive constraints (larger yellow space) superimposed to the region determined from the inclusive mode (smaller inner space). The top panel refers to $\mathcal{B}(B^- \rightarrow D^0 \mu^- \bar{\nu}_\mu)$, the middle one to $\frac{d\Gamma}{dw}(B^- \rightarrow D^{*0} \mu^- \bar{\nu}_\mu)$. In each case there is an overlap between the regions selected from the inclusive and the exclusive mode, and a region exists where all constraints are fulfilled, as shown in the bottom panel. The exclusive data slightly reduce the upper bound on $|V_{cb}|$ and produce a lower bound. This is shown in Fig. 3, where we depict the projections in the $(\text{Re}(\epsilon_T^\mu), \text{Im}(\epsilon_T^\mu))$ plane of the parameter space corresponding to the extreme values of $|V_{cb}|$, i.e., the values for which the parameter regions determined through the various constraints do not overlap. After the application of all constraints, the range for $|V_{cb}|$ is $|V_{cb}| \in [0.0343, 0.0421]$. The range for ϵ_T^μ depends on $|V_{cb}|$, and the largest allowed value is $|\epsilon_T^\mu| \simeq 0.2$. The symmetry axes of the two regions of parameters bound through the inclusive and the exclusive constraints do not intersect the $(\text{Re}(\epsilon_T^\mu), \text{Im}(\epsilon_T^\mu))$ plane at the origin and do not coincide. This is due to the lepton mass effect and to the interference term in the rates. The NP contribution affects in a different way the inclusive and the two exclusive B channels.

For the electron mode the analysis is repeated, with changes in the results since the electron mass is tiny. We show in Fig. 4 the parameter space selected from the inclusive and the exclusive D mode (top panel), and from the inclusive and the D^* mode (middle panel). The three constraints are fulfilled for parameters in the region depicted in the bottom panel. The upper bound on $|V_{cb}|$ selected from the inclusive analysis, $|V_{cb}| \leq 0.04273$, is not modified by the exclusive constraints, while a lower bound is found, as it can be inferred from Fig. 5. The result from the electron modes is the range

$$|V_{cb}| \in [0.036, 0.0427] \text{ and the bound } |\epsilon_T^e| \leq 0.17.$$

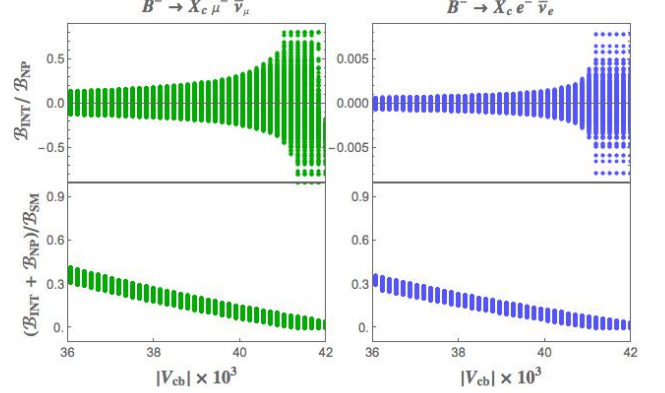


FIG. 6: Relative size of the NP contributions to $\mathcal{B}(B^- \rightarrow X_c \mu^- \bar{\nu}_\mu)$ (left) and $\mathcal{B}(B^- \rightarrow X_c e^- \bar{\nu}_e)$ (right).

The conclusion is that there are sets of three parameters fulfilling all the constraints. For $|V_{cb}|$ this happens in the range $|V_{cb}| \in [0.036, 0.042]$. Although this implies a sizeable uncertainty on $|V_{cb}|$, the analysis shows that a non SM contribution, as the one considered here, can reconcile its inclusive and exclusive determination.

To study the role of the NP contribution and of the interference between SM and NP, in the inclusive channel we separately integrate the three terms in (4). Denoting the resulting quantities, divided by the full decay width, as \mathcal{B}_{SM} , \mathcal{B}_{NP} , \mathcal{B}_{INT} , we consider the ratios $\frac{\mathcal{B}_{INT}}{\mathcal{B}_{NP}}$ and $\frac{\mathcal{B}_{INT} + \mathcal{B}_{NP}}{\mathcal{B}_{SM}}$, and display in Fig. 6 the results obtained varying $(\text{Re}(\epsilon_T^\ell), \text{Im}(\epsilon_T^\ell), |V_{cb}|)$ in their allowed region. The interference term \mathcal{B}_{INT} is sizable with respect to \mathcal{B}_{NP} for muon. Moreover, for both μ and e , the NP contribution $\mathcal{B}_{INT} + \mathcal{B}_{NP}$ becomes negligible with respect to SM for large $|V_{cb}|$, while it is sizable for smaller values of $|V_{cb}|$. The analogous quantities for the exclusive $B^- \rightarrow D^0 \ell^- \bar{\nu}_\ell$ modes are in Fig. 7. The interference term

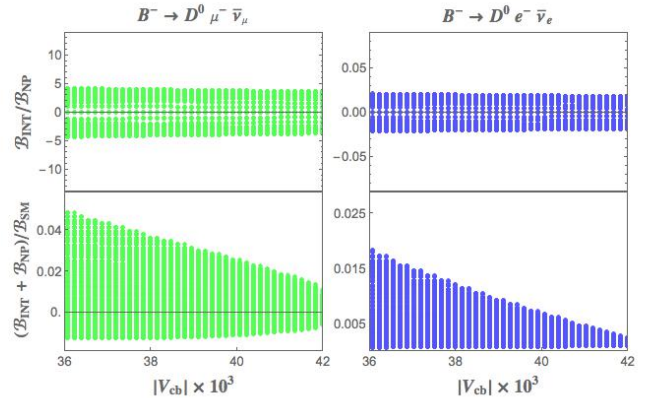


FIG. 7: Relative size of the NP contributions to $\mathcal{B}(B^- \rightarrow D^0 \mu^- \bar{\nu}_\mu)$ (left) and $\mathcal{B}(B^- \rightarrow D^0 e^- \bar{\nu}_e)$ (right).

is again non negligible for muon. The impact of the new operator is larger in the inclusive mode than in D . The changes due to NP are different in the different channels, showing how an extra term in the effective Hamiltonian could be at the origin of the $|V_{cb}|$ anomaly.

Considering the ranges determined for the couplings ϵ_T^e and ϵ_T^μ , it is possible that a new physics contribution of the kind considered here still satisfies $e - \mu$ universality. Improved measurements and reduced theoretical uncertainties, mainly in the hadronic form factors, are needed to shed light on this point.

Conclusions. We have given an example of a possible

mechanism at the origin of the tension in $|V_{cb}|$ measurements. This mechanism can be tested in other processes, namely B_s and Λ_b semileptonic decays, for which enough information is not yet available. A determination of $|V_{cb}|$ from the purely leptonic $B_c \rightarrow \ell \bar{\nu}$ decay is interesting, since in that case the tensor operator in Eq. (3) does not contribute. As for the experimental investigations, the importance of separate measurements of the muon and electron modes cannot be overemphasized.

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